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Hyperon Masses in Nuclear Matter *

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Abstract

We analyze hyperon and nucleon mass shifts in nuclear matter using chiral perturbation theory. Expressions for the mass shifts that include strong interaction effects at leading order in the density are derived. Corrections to our results are suppressed by powers of the Fermi momentum divided by either the chiral symmetry breaking scale or the nucleon mass. Our work is relevant for neutron stars and for large hypernuclei.

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In nuclear matter the masses of hyperons and nucleons are shifted from their vacuum values by strong interaction effects. Understanding these mass shifts is important for neutron stars [1–7] and for large hypernuclei (see for example [8,9]). Most approaches to this problem rely on the use of phenomenological models [1,2,4–6,9–15] since the relevant strong interaction effects involve QCD in the nonperturbative regime. In this letter we will consider mass shifts for the baryon octet using a model independent approach to baryon-baryon interactions based on chiral perturbation theory. The price to be paid for model independence is that an expansion in momentum is necessary and that our expressions for the mass shifts involve parameters in the heavy baryon chiral Lagrangian that have not yet been determined from experiment. However, there are mass relations that hold independent of these parameters in analogy with the Gell-Mann–Okubo mass formula for SU(3) breaking.

We study changes in the masses of the baryon octet $(n, p, \Lambda, \Sigma, \Xi)$ that arise from their strong interactions with a degenerate Fermi gas of neutrons at density $d^{(n)} = (p_F^{(n)})^3/3\pi^2$ and protons at density $d^{(p)} = (p_F^{(p)})^3/3\pi^2$ (i.e., nuclear matter). These interactions are determined by a heavy-baryon chiral Lagrangian that is consistent with the $SU(3)_L \times SU(3)_R$ chiral symmetry of QCD. An expansion in derivatives is appropriate for Fermi momenta $p_F^{(n)}$ and $p_F^{(p)}$ that are small compared with the chiral symmetry breaking scale and the baryon masses.

In the rest frame of background nuclear matter the energy of a baryon and its three-momentum are related (for $|\vec{p}| << M_B$) by $E^{\rm NM} = M_B + \Delta M_B + |\vec{p}|^2/2M_B + \eta|\vec{p}|^2 + \ldots$ which is to be compared to the free space relation $E^{\rm vac} = M_B + |\vec{p}|^2/2M_B + \ldots$ where the dots denote terms higher order in the momentum expansion. We can combine the effects of the medium with the free space mass, M_B, by defining M_B^* and $\tilde{M}_B$, with $\tilde{M}_B = M_B + \Delta M_B$ and $M_B^* = M_B(1+2\eta M_B)^{-1}$ allowing us to write $E^{\rm NM} = \tilde{M}_B + |\vec{p}|^2/2M_B^* + \ldots$. The quantities ΔM_B and η arise from interactions with the background medium and in general $M_B^* \neq \tilde{M}_B$. We compute the lowest order contribution to the momentum independent mass shift ΔM_B for the octet baryons using chiral perturbation theory.

The pseudo-Goldstone boson fields (π, K, η) can be written as a 3×3 special unitary matrix

$$\Sigma = \exp \frac{2i\mathbf{\Pi}}{f} \quad , \tag{1}$$

where

$$\mathbf{\Pi} = \begin{bmatrix} \pi^0/\sqrt{2} + \eta/\sqrt{6} & \pi^+ & K^+ \\ \pi^- & -\pi^0/\sqrt{2} + \eta/\sqrt{6} & K^0 \\ K^- & \bar{K}^0 & -\sqrt{\frac{2}{3}}\eta \end{bmatrix} \quad . \tag{2}$$

Under chiral $SU(3)_L \times SU(3)_R$ symmetry, $\Sigma \rightarrow L\Sigma R^\dagger$ where $L \epsilon SU(3)_L$ and $R \epsilon SU(3)_R$. At leading order in chiral perturbation theory f can be identified with the pion decay constant ($f_\pi \simeq 132 MeV$). When describing the interactions of the pseudo-Goldstone bosons with other fields it is convenient to introduce

$$\xi = \exp \frac{i\mathbf{\Pi}}{f} = \sqrt{\Sigma} \quad . \tag{3}$$

Under chiral $SU(3)_L \times SU(3)_R$ symmetry

$$\xi \to L\xi U^\dagger = U\xi R^\dagger \quad , \tag{4}$$

where in general U is a complicated function of L, R and the meson fields $\mathbf{\Pi}$. For transformations $V = L = R$ in the unbroken $SU(3)_V$ subgroup $U = V$.

The baryon fields are introduced as a 3×3 octet matrix

$$B = \begin{bmatrix} \Sigma^0/\sqrt{2} + \Lambda/\sqrt{6} & \Sigma^+ & p \\ \Sigma^- & -\Sigma^0/\sqrt{2} + \Lambda/\sqrt{6} & n \\ \Xi^- & \Xi^0 & -\sqrt{\frac{2}{3}}\Lambda \end{bmatrix} \quad , \tag{5}$$

that transforms under chiral $SU(3)_L \times SU(3)_R$ as

$$B \to UBU^\dagger \quad . \tag{6}$$

We construct a chiral Lagrangian by treating the baryons as heavy static 2-component fields. The chiral Lagrangian is written as

$$\mathcal{L} = \mathcal{L}^{(1)} + \mathcal{L}^{(2)} + \ldots \quad , \tag{7}$$

where $\mathcal{L}^{(j)}$ contains $2j$ baryon fields.

$\mathcal{L}^{(1)}$ is the familiar heavy baryon chiral Lagrangian that gives the interactions of the baryon octet with the pseudo-Goldstone bosons. At leading order in chiral perturbation theory

$$\begin{aligned} \mathcal{L}^{(1)} = & \, TrB_j^\dagger i\partial_0 B_j + iTrB_j^\dagger[V_0, B_j] \\ & - DTrB_j^\dagger \vec{\sigma}_{jk}\{\vec{A}, B_k\} - FTrB_j^\dagger \vec{\sigma}_{jk}[\vec{A}, B_k] \quad , \end{aligned} \tag{8}$$

with the repeated spin indices j and k summed over 1,2 and the vector and axial-vector chiral fields are

$$V_\mu = \frac{1}{2}(\xi^\dagger \partial_\mu \xi + \xi \partial_\mu \xi^\dagger) \quad , \tag{9}$$

$$A_\mu = \frac{i}{2}(\xi^\dagger \partial_\mu \xi - \xi \partial_\mu \xi^\dagger) \quad . \tag{10}$$

Nuclear β decay and semileptonic hyperon decay give $F \simeq 0.44$ and $D \simeq 0.81$ at tree-level. At higher order in chiral perturbation theory there are corrections to Eq. (8) coming from terms with more derivatives and terms with insertions of the light quark mass matrix.

Interactions between baryons mediated by pseudo-Goldstone boson exchange give rise to hyperon mass shifts. At the same order in chiral perturbation theory (i.e., order p_F^3) terms in $\mathcal{L}$ with four baryon fields (and no derivatives) also play a role. They are given by

$$\begin{aligned} \mathcal{L}^{(2)} = & - \frac{c_1}{f^2} Tr(B_i^\dagger B_i B_j^\dagger B_j) - \frac{c_2}{f^2} Tr(B_i^\dagger B_j B_j^\dagger B_i) \\ & - \frac{c_3}{f^2} Tr(B_i^\dagger B_j^\dagger B_i B_j) - \frac{c_4}{f^2} Tr(B_i^\dagger B_j^\dagger B_j B_i) \\ & - \frac{c_5}{f^2} Tr(B_i^\dagger B_i) Tr(B_j^\dagger B_j) - \frac{c_6}{f^2} Tr(B_i^\dagger B_j) Tr(B_j^\dagger B_i) \quad . \end{aligned} \tag{11}$$

Factors of $1/f^2$ appear in Eq. (11) so that the coefficients c_j are dimensionless. At higher order in chiral perturbation theory there are corrections to Eq. (11) coming from terms with derivatives and terms with insertions of the light quark mass matrix. There are also contributions to the baryon mass shifts from operators in $\mathcal{L}^{(j)}$, $j > 2$, involving more baryon fields, e.g. $Tr(B_i^\dagger B_i B_j^\dagger B_j B_k^\dagger B_k)$, but these are suppressed by additional powers of $d^{(n,p)}/(4\pi f^3)$ over the contribution from operators appearing in Eq. (11) . Note there is no term involving $Tr(B_i^\dagger B_j^\dagger)\ Tr(B_i B_j)$ in Eq. (11) as it can be eliminated using a Cayley-Hamilton operator identity:

$$
\begin{aligned}
&-Tr(B_i^\dagger B_j^\dagger B_i B_j) + Tr(B_i^\dagger B_j^\dagger B_j B_i) - \frac{1}{2} Tr(B_i^\dagger B_j B_j^\dagger B_i) + \frac{1}{2} Tr(B_i^\dagger B_i B_j^\dagger B_j) \\
&= \frac{1}{2} Tr(B_i^\dagger B_i) Tr(B_j^\dagger B_j) - \frac{1}{2} Tr(B_i^\dagger B_j) Tr(B_j^\dagger B_i) - \frac{1}{2} Tr(B_i^\dagger B_j^\dagger) Tr(B_i B_j) \quad , \qquad (12)
\end{aligned}
$$

It is easy to understand from spin-flavour considerations why there are only six independent four-baryon operators. The $SU(3)_V$ decompostion of the product of two baryon octets is $\mathbf{8} \otimes \mathbf{8} = \mathbf{27} \oplus \mathbf{10} \oplus \overline{\mathbf{10}} \oplus \mathbf{8}_A \oplus \mathbf{8}_S \oplus \mathbf{1}$ of which the $\mathbf{27}, \mathbf{8}_S, \mathbf{1}$ are symmetric and the $\mathbf{10}, \overline{\mathbf{10}}, \mathbf{8}_A$ are antisymmetric under baryon interchange. Fermi statistics require that the baryons in the $\mathbf{27}, \mathbf{8}_S, \mathbf{1}$ have spin $S = 0$ and those in the $\mathbf{10}, \overline{\mathbf{10}}, \mathbf{8}_A$ have $S = 1$ as the operators under consideration give rise to S-wave interactions only. Hence, there are only 6 independent four-baryon operators that exist at this order in chiral perturbation theory. It is these six operators (and pseudo-Goldstone boson exchange) that give the leading contribution to S-wave baryon-baryon scattering, such as $\Lambda N \to \Lambda N$, in chiral perturbation theory. An $SU(3)_V$ analysis of hyperon-nucleon scattering can be found in [16].

The four-baryon operators in Eq. (11) include effects from the repulsive core of the baryon-baryon potential. For example, a repulsive spherical well of infinite height and width r_0 gives rise to coefficients, c_j, of order $4\pi f^2 r_0/M_B$.

Using Eq. (8) and Eq. (11) we find the following one loop expressions for the baryon octet mass shifts in nuclear matter:

$$
\begin{aligned}
\Delta M_n &= \frac{d^{(n)}}{f^2}\left\{ -\frac{(F+D)^2}{4}\mathcal{I}(m_\pi/p_F^{(n)}) - \frac{(3F-D)^2}{12}\mathcal{I}(m_\eta/p_F^{(n)}) + c_1 - c_2 + c_5 - c_6 \right\} \\
&\quad + \frac{d^{(p)}}{f^2}\left\{ -\frac{(F+D)^2}{2}\mathcal{I}(m_\pi/p_F^{(p)}) + 2c_1 + c_2 + 2c_5 + c_6 \right\} \\
\Delta M_p &= \frac{d^{(n)}}{f^2}\left\{ -\frac{(F+D)^2}{2}\mathcal{I}(m_\pi/p_F^{(n)}) + 2c_1 + c_2 + 2c_5 + c_6 \right\} \\
&\quad + \frac{d^{(p)}}{f^2}\left\{ -\frac{(F+D)^2}{4}\mathcal{I}(m_\pi/p_F^{(p)}) - \frac{(3F-D)^2}{12}\mathcal{I}(m_\eta/p_F^{(p)}) + c_1 - c_2 + c_5 - c_6 \right\} \\
\Delta M_{\Sigma^+} &= \frac{d^{(n)}}{f^2}\left\{ \frac{c_3}{2} + c_4 + 2c_5 + c_6 \right\} \\
&\quad + \frac{d^{(p)}}{f^2}\left\{ -\frac{(D-F)^2}{2}\mathcal{I}(m_K/p_F^{(p)}) - c_1 - 2c_2 + 2c_5 + c_6 \right\} \\
\Delta M_{\Sigma^-} &= \frac{d^{(n)}}{f^2}\left\{ -\frac{(D-F)^2}{2}\mathcal{I}(m_K/p_F^{(n)}) - c_1 - 2c_2 + 2c_5 + c_6 \right\}
\end{aligned}
$$

$$+\frac{d^{(p)}}{f^2}\left\{\frac{c_3}{2}+c_4+2\,c_5+c_6\right\}$$

$$\Delta M_{\Xi^0}=\frac{d^{(n)}}{f^2}\left\{c_3+2\,c_4+2\,c_5+c_6\right\}+\frac{d^{(p)}}{f^2}\left\{\frac{1}{2}c_3+c_4+2\,c_5+c_6\right\}$$

$$\Delta M_{\Xi^-}=\frac{d^{(n)}}{f^2}\left\{\frac{1}{2}c_3+c_4+2\,c_5+c_6\right\}+\frac{d^{(p)}}{f^2}\left\{c_3+2\,c_4+2\,c_5+c_6\right\}\quad . \tag{13}$$

For the Λ and Σ^0 there is a 2×2 mass matrix with entries

$$\Delta M_{\Lambda\Lambda}=\frac{d^{(n)}}{f^2}\left\{-\frac{1}{12}(D+3F)^2\mathcal{I}(m_K/p_F^{(n)})+\frac{7}{6}c_1+\frac{1}{3}c_2+\frac{13}{12}c_3+\frac{7}{6}c_4+2c_5+c_6\right\}$$
$$+\frac{d^{(p)}}{f^2}\left\{-\frac{1}{12}(D+3F)^2\mathcal{I}(m_K/p_F^{(p)})+\frac{7}{6}c_1+\frac{1}{3}c_2+\frac{13}{12}c_3+\frac{7}{6}c_4+2c_5+c_6\right\}$$
$$\Delta M_{\Sigma^0\Sigma^0}=\frac{d^{(n)}}{f^2}\left\{-\frac{(D-F)^2}{4}\mathcal{I}(m_K/p_F^{(n)})-\frac{1}{2}c_1-c_2+\frac{1}{4}c_3+\frac{1}{2}c_4+2c_5+c_6\right\}$$
$$+\frac{d^{(p)}}{f^2}\left\{-\frac{(D-F)^2}{4}\mathcal{I}(m_K/p_F^{(p)})-\frac{1}{2}c_1-c_2+\frac{1}{4}c_3+\frac{1}{2}c_4+2c_5+c_6\right\}$$
$$\Delta M_{\Lambda\Sigma}=\Delta M_{\Sigma\Lambda}=\frac{d^{(n)}}{f^2}\left\{\frac{(D+3F)(F-D)}{4\sqrt{3}}\mathcal{I}(m_K/p_F^{(n)})+\frac{c_1}{2\sqrt{3}}+\frac{c_2}{\sqrt{3}}-\frac{5}{4\sqrt{3}}c_3-\frac{1}{\sqrt{3}}c_4\right\}$$
$$+\frac{d^{(p)}}{f^2}\left\{\frac{(D+3F)(D-F)}{4\sqrt{3}}\mathcal{I}(m_K/p_F^{(p)})-\frac{c_1}{2\sqrt{3}}-\frac{c_2}{\sqrt{3}}+\frac{5}{4\sqrt{3}}c_3+\frac{1}{\sqrt{3}}c_4\right\}\quad . \tag{14}$$

In Eq. (13) and Eq. (14) the function $\mathcal{I}(x)$ is defined by

$$\mathcal{I}(x)=1-3x^2+3x^3 Arctan(1/x)\quad . \tag{15}$$

Naively, corrections to Eq. (13) and Eq. (14) are suppressed by powers of the Fermi momentum divided by the chiral symmetry breaking scale (or the nucleon mass) and by powers of the light quark masses divided by the chiral symmetry breaking scale (or the nucleon mass). This is essentially a consequence of dimensional analysis as Feynman diagrams with more than one loop constructed from the vertices of the Lagrange densities in Eq. (8) and Eq. (11) contain more powers of $1/f$ which must be compensated by additional factors of p_F or light quark masses. Corrections to the Lagrange densities in Eq. (8) and Eq. (11) from operators with more derivatives (or insertions of the light quark mass matrix) have coefficients with more powers of $1/f$ which again must be compensated by additional factors of p_F or light quark masses. However, this power counting is not quite correct. Feynman diagrams with more baryon propagators than loops have infrared divergent loop integrations over energy. Furthermore, insertions of the kinetic energy, $\mathcal{L}_{KE}=TrB_i^\dagger\nabla^2 B_i/2M_B$, increase the number of baryon propagators but not the number of loops (this correction to the Lagrange density in Eq. (8) cannot be treated as perturbation). The cure for these problems is to sum insertions of the kinetic energy changing the baryon propagator to

$$i\left[\frac{\Theta(p_F-|\vec{p}|)}{p^0-\vec{p}^2/2M_B-i\epsilon}+\frac{\Theta(|\vec{p}|-p_F)}{p^0-\vec{p}^2/2M_B+i\epsilon}\right]\quad . \tag{16}$$

Then infrared divergent p^0 integrations are cut off by the kinetic energy which produces factors of M_B in the numerator. At l-loops one finds corrections suppressed by a factor of order $(M_B p_F/(4\pi f^2))^{\ell-1}$. As a consequence, the leading correction to our results arise from two loop diagrams and are suppressed by only $\sim M_B p_F/(4\pi f^2)$ (instead of $\sim p_F^2/(4\pi f)^2$). At nuclear density ($p_F \sim 350$ MeV/c) this factor is not small. Therefore the leading order results derived here may receive significant corrections from higher orders in chiral perturbation theory.

We adopt power counting rules where M_B is considered to be of the same order as the chiral symmetry breaking scale $\sim 4\pi f$ and two factors of p_F are considered to be of the same order as a light quark mass (Recall that the squares of the pseudo-Goldstone boson masses are proportional to light quark masses). Consequently in Eq. (13) and Eq. (14) the full dependence of the baryon mass shifts on (p_F/μ) where $\mu = m_K, m_\pi$ or m_η is kept.

The baryon mass shifts given in Eq. (13) and Eq. (14) are our principal results. If the nuclear matter is not isospin symmetric (i.e., $d^{(n)} \neq d^{(p)}$) there is $\Lambda - \Sigma^0$ mixing (Eq. (14)). More importantly, the eight octet masses are expressed in terms of six coefficients c_j so there are mass relations independent of the coefficients of the four baryon operators. We find that

$$2\Delta M_{\Sigma^0\Sigma^0} - \Delta M_{\Sigma^+} - \Delta M_{\Sigma^-} = 0 \quad , \tag{17}$$

which is independent of the composition of the nuclear matter, it is true for an arbitrary proton to neutron ratio. A second mass relation is given in terms of the mesonic contribution

$$d^{(n)}\left(\Delta M_{\Sigma^+} - \Delta M_{\Xi^-}\right) + d^{(p)}\left(\Delta M_{\Xi^0} - \Delta M_{\Sigma^-}\right) = \frac{(D-F)^2}{2f^2} d^{(n)} d^{(p)} \left(\mathcal{I}(m_K/p_F^{(n)}) - \mathcal{I}(m_K/p_F^{(p)})\right) \quad . \tag{18}$$

We see that in the extreme case of neutron matter ($d^{(p)} = 0$) this relation becomes $\Delta M_{\Sigma^+} = \Delta M_{\Xi^-}$. There is also a relation involving the $\Sigma^0\Lambda$ mixing term

$$\begin{aligned}
& f^2 d^{(n)}\left(-3\Delta M_{\Lambda\Lambda} + 2\Delta M_p - 2\sqrt{3}\Delta M_{\Lambda\Sigma} - 2\Delta M_{\Sigma^-} + 3\Delta M_{\Sigma^0\Sigma^0}\right) \\
+ & f^2 d^{(p)}\left(3\Delta M_{\Lambda\Lambda} - 2\Delta M_n - 2\sqrt{3}\Delta M_{\Lambda\Sigma} - 2\Delta M_{\Sigma^-} + \Delta M_{\Sigma^0\Sigma^0}\right) \\
& = -(D+F)^2\left[{d^{(n)}}^2 + {d^{(p)}}^2 - \frac{1}{2} d^{(n)} d^{(p)}\right]\left(\mathcal{I}(m_\pi/p_F^{(n)}) - \mathcal{I}(m_\pi/p_F^{(p)})\right) \\
& \quad + (D+F)(D-3F)\, d^{(n)} d^{(p)}\left(\mathcal{I}(m_K/p_F^{(n)}) - \mathcal{I}(m_K/p_F^{(p)})\right) \\
& \quad + (D+F)^2\left({d^{(n)}}^2 \mathcal{I}(m_K/p_F^{(n)}) - {d^{(p)}}^2 \mathcal{I}(m_K/p_F^{(p)})\right) \\
& \quad + \frac{1}{6}(D-3F)^2\, d^{(n)} d^{(p)}\left(\mathcal{I}(m_\eta/p_F^{(n)}) - \mathcal{I}(m_\eta/p_F^{(p)})\right)
\end{aligned} \tag{19}$$

When $d^{(p)} = d^{(n)}$ the background nuclear matter is isospin symmetric leading to baryons in the same isomultiplet receiving the same mass shift and $\Sigma^0\Lambda$ mixing vanishing.

Weinberg [17,18] has extracted values for $c_2 + c_6 = f^2 C_T$ and $2c_1 + c_2 + 2c_5 + c_6 = f^2 C_S$ from the measured S-wave nucleon scattering lengths . He found $C_T \approx (1/45\mathrm{MeV})^2$ and $C_S \approx -(1/88\mathrm{MeV})^2$ which give very large nucleon mass shifts at nuclear density [19]. If Eq. (13) and Eq. (14) are to be relevant at nuclear density the coefficients $C_{S,T}$ must be smaller than these values. It may be possible to extract other combinations of the coefficients c_i by measuring the S-wave scattering lengths for hyperon nucleon scattering. For hyperon-proton strong interactions the relevant spin singlet $a_{S=0}$ and spin triplet $a_{S=1}$ elastic scattering lengths are

$$\begin{aligned}
|a_{S=0}(\Lambda p \to \Lambda p)| &= \frac{1}{2\pi f^2}\frac{M_p M_\Lambda}{M_p + M_\Lambda}|\frac{5}{3}c_1 - \frac{5}{3}c_2 - \frac{1}{6}c_3 + \frac{1}{6}c_4 + 2c_5 - 2c_6| \\
|a_{S=1}(\Lambda p \to \Lambda p)| &= \frac{1}{2\pi f^2}\frac{M_p M_\Lambda}{M_p + M_\Lambda}|c_1 + c_2 + \frac{3}{2}c_3 + \frac{3}{2}c_4 + 2c_5 + 2c_6| \\
|a_{S=0}(\Sigma^+ p \to \Sigma^+ p)| &= \frac{1}{2\pi f^2}\frac{M_p M_\Sigma}{M_p + M_\Sigma}2|c_1 - c_2 + c_5 - c_6| \\
|a_{S=1}(\Sigma^+ p \to \Sigma^+ p)| &= \frac{1}{2\pi f^2}\frac{M_p M_\Sigma}{M_p + M_\Sigma}2| - c_1 - c_2 + c_5 + c_6| \\
|a_{S=0}(\Sigma^- p \to \Sigma^- p)| &= \frac{1}{2\pi f^2}\frac{M_p M_\Sigma}{M_p + M_\Sigma}| - c_3 + c_4 + 2c_5 - 2c_6| \\
|a_{S=1}(\Sigma^- p \to \Sigma^- p)| &= \frac{1}{2\pi f^2}\frac{M_p M_\Sigma}{M_p + M_\Sigma}|c_3 + c_4 + 2c_5 + 2c_6| \\
|a_{S=0}(\Xi^- p \to \Xi^- p)| &= \frac{1}{2\pi f^2}\frac{M_p M_\Xi}{M_p + M_\Xi}2| - c_3 + c_4 + c_5 - c_6| \\
|a_{S=1}(\Xi^- p \to \Xi^- p)| &= \frac{1}{2\pi f^2}\frac{M_p M_\Xi}{M_p + M_\Xi}2|c_3 + c_4 + c_5 + c_6| \\
|a_{S=0}(\Xi^0 p \to \Xi^0 p)| &= \frac{1}{2\pi f^2}\frac{M_p M_\Xi}{M_p + M_\Xi}|c_3 - c_4 - 2c_5 + 2c_6| \\
|a_{S=1}(\Xi^0 p \to \Xi^0 p)| &= \frac{1}{2\pi f^2}\frac{M_p M_\Xi}{M_p + M_\Xi}|c_3 + c_4 + 2c_5 + 2c_6| \qquad (20)
\end{aligned}$$

Because of the derivative coupling in Eq. (8) tree-level pseudo-Goldstone boson exchange doesn't contribute to the scattering lengths.

The baryon mass shifts given in Eq. (13) and Eq. (14) may be of use for large hypernuclei and for neutron stars. Assuming a large nucleus can be modeled by a degenerate Fermi gas replacing one of its neutrons with a Λ results in a hypernucleus with a mass shift that is partly determined by $\Delta M_\Lambda - \Delta M_n$. Study of the S-wave hyperon-nucleon scattering lengths may eventually lead to a determination of the coefficients c_j which can then be used to predict hyperon mass shifts in neutron stars (where $d^{(p)} \ll d^{(n)}$) and in large hypernuclei.

The utility of our results depends crucially upon the smallness of higher orders in chiral perturbation theory at nuclear density. A computation at next order in chiral perturbation theory may help resolve this issue. We hope to present results on this in a future publication.

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